Determination of the angle γ from nonleptonic $B_c \to D_s D^0$ decays

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February 1, 2008

Abstract

We note that the two body nonleptonic pure tree decays $B_c^\pm \to D_s^\pm D^0(\bar{D}^0)$ and the corresponding vector-vector modes $B_c^\pm \to D_s^{*\pm} D^{*0}(\bar{D}^{*0})$ are well suited to extract the weak phase γ of the unitarity triangle. The CP violating phase γ can be determined cleanly as these decay modes are free from the penguin pollutions.

PACS Nos. : 11.30 Er, 12.15 Hh, 13.25 Hw

1 Introduction

Despite many attempts, CP violation still remains one of the most outstanding problems in particle physics [1, 2]. The standard model (SM) with three generations provides a simple description of this phenomenon through the complex Cabibbo-Kabayashi-Maskawa matrix [3]. Decays of B mesons provide rich ground for investigating CP violation [4, 5]. They allow stringent tests both for the SM and for studies of new sources of this effect. Within the SM, the CP violation is often characterized by the so-called unitarity triangle [6]. Detection of CP violation and the accurate determination of the unitarity triangle are the major goals of experimental B physics [7]. Decisive information about the origin of CP violation in the flavor sector can be obtained if the three angles $\alpha(\equiv \phi_2)$, $\beta(\equiv \phi_1)$ and $\gamma(\equiv \phi_3)$ can be independently measured [8]. The sum of these three angles must be equal to 180° if the SM with three generations is the model for the CP violation. These angles of the unitarity triangle can be loosely bounded from various low energy phenomenology. The angle $\beta(\equiv \phi_1)$ can be measured from the gold plated mode $B_d \to J/\psi K_s$ without any hadronic uncertainty. The angle $\alpha(\equiv \phi_2)$ can be measured from $B \to \pi\pi$ mode but there is some penguin contamination. Still one can hope to perform the isospin analysis and remove the penguin contribution thereby extracting

the angle $\alpha(\phi_2)$ with reasonable accuracy [9]. The most difficult to measure is the angle $\gamma(\equiv \phi_3)$. There have been a lot of suggestions and discussions about how to measure this quantity at B factories [10, 11]. In Ref. [10] the authors proposed to extract γ using the independent measurements of $B \to D^0 K$, $B \to \bar{D}^0 K$ and $B \to D_{CP} K$. However, for the charged B meson decay mode $A(B^- \to \bar{D}^0 K^-)$ is difficult to measure experimentally. The reason is that the final $\bar{D}^0$ meson should be identified using $\bar{D}^0 \to K^+ \pi^-$ but it is difficult to distinguish it from doubly Cabibbo suppressed $D^0 \to K^+ \pi^-$. There are various methods to overcome these difficulties. In Ref. [12] Atwood et al used different final states into which the neutral D meson decays, to extract information on γ. In Ref. [13] Gronau proposed that the angle γ can be determined by using the color allowed decay modes $B^- \to D^0 K^-$, $B^- \to D_{CP} K^-$ and their charge conjugation modes. In Ref. [14] a new method, using the isospin relations, is suggested to extract γ by exploiting the decay modes $B \to D K^{(*)}$ that are not Cabibbo suppressed. Falk and Petrov [15] recently proposed a new method for measuring γ using the partial rates for CP-tagged B_s decays. The angle γ can also be measured using the $SU(3)$ relations between $B \to \pi K, \pi\pi$ decay amplitudes [16]. Although this approach is not theoretically clean in contrast to the $B \to KD$ strategies using pure tree decays - it is more promising from experimental point of view. In Ref. [17] it is proposed that γ can be determined using the $B \to D^* V$ $(V = K^*, \rho)$ modes.

The decays of B_c meson ($\bar{c}b$ and $\bar{b}c$ bound states) seem to be another valuable window for probing the origin of CP violation. Since large number of B_c meson is expected to be produced at hadronic colliders like LHC or Tevatron, it is therefore interesting to examine the features of CP violation in B_c mesons. In this paper we would like to discuss about the determination of angle γ from the pure tree nonleptonic decay modes $B_c^\pm \to D_s^\pm \{D^0, \bar{D}^0, D_{CP}^\pm\}$ and the vector vector modes $B_c^\pm \to D_s^{*\pm} D^{*0}(\bar{D}^{*0})$. In Ref. [10] the analogous pseudoscalar decay modes i.e., $B^\pm \to D^0(\bar{D}^0) K^\pm$, are considered for the determination of the angle γ. However, the corresponding B_c counterpart has some additional advantages. The isospin analysis done by Deshpande and Dib [18] for the former case implies that as D and K are isospin 1/2 objects, a DK final state can be either $I = 0$ or 1. Since strong interactions conserve isospin, in general there will be two strong rescattering phases δ_0 and δ_1 , one for each final state of given isospin. However, in case of $B_c^\pm \to D_s^\pm D^0(\bar{D}^0)$ modes, the initial B_c state is an isosinglet and the final states $D_s^\pm D^0(\bar{D}^0)$ are isodoublets. Hence these processes are described by the effective Hamiltonian of $|\Delta I| = 1/2$. Since isospin is conserved in strong interactions in general one can expect same strong phases for all the final states. Hence the uncertainties due to the presence of the notorious strong phases can be eliminated without any additional assumptions and γ can be determined cleanly. But in actual practice there could be some amount of strong phase difference between these amplitudes from resonance effects [19], which we are not taking into account

in this analysis. It has been shown recently by Fleischer and Wyler [20] that the B_c counterpart of $B^\pm \to K^\pm D$ triangle approach be well suited to extract the angle γ as both the amplitudes $B_c^+ \to D_s^+ D^0$ and $B_c^+ \to D_s^+ \bar{D}^0$ are of the same order of magnitude.

The paper is organised as follows. We present the method for the determination of the angle γ from the decay mode $B_c^\pm \to D_s^\pm D^0(\bar{D}^0)$ in section II and from the vector vector modes $B_c^\pm \to D_s^{*\pm} D^{*0}(\bar{D}^{*0})$ in section III. Section IV contains our conclusion.

2 γ from $B_c^\pm \to D_s^\pm D^0(\bar{D}^0)$

The effective Hamiltonians for the decay modes $B_c^- \to D_s^- D^0$ and $B_c^- \to D_s^- \bar{D}^0$, described by the quark level transitions $b \to c\bar{u}s$ and $b \to u\bar{c}s$ respectively are given as

$$\begin{aligned}
\mathcal{H}_{eff}(b \to c\bar{u}s) &= \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* \left[C_1(m_b)(\bar{s}u)(\bar{c}b) + C_2(m_b)(\bar{c}u)(\bar{s}b)\right] \\
\mathcal{H}_{eff}(b \to u\bar{c}s) &= \frac{G_F}{\sqrt{2}} V_{ub} V_{cs}^* \left[C_1(m_b)(\bar{s}c)(\bar{u}b) + C_2(m_b)(\bar{u}c)(\bar{s}b)\right] \qquad (1)
\end{aligned}$$

where C_1 and C_2 are the Wilson coefficients with values evaluated at the b-quark mass scale as [21]

$$C_1(m_b) = 1.13, \quad \text{and} \quad C_2(m_b) = -0.29\ , \qquad (2)$$

$(\bar{c}b) = \bar{c}\gamma_\mu(1-\gamma_5)b$ etc. are the usual $(V-A)$ color singlet quark currents. The hadronic matrix elements of the four quark current operators i.e., $\langle D_s D|\mathcal{H}_{eff}|B_c\rangle$ are very difficult to evaluate from the first principle of QCD. The usual way to evaluate these matrix elements for nonleptonic B-decays is to assume some approximation. Here we use the factorization approximation which factorizes each four quark matrix element into a product of two elements. In the naive factorization hypothesis only the factorizable contributions are considered. However, the nonfactorizable amplitudes which cannot be calculated in the naive factorization approach are important for understanding the data. So the generalized factorization approach is assumed where these nonfactorizable contributions are incorporated in a phenomenological way: they are lumped into the coefficients $a_1 = C_1 + C_2/N_c$ and $a_2 = C_2 + C_1/N_c$, so that now the effective coefficients a_1^{eff} and a_2^{eff} are treated as free parameters and their values can be extracted from the experimental data. From now onwards we shall denote a_1^{eff} and a_2^{eff} by simply a_1 and a_2. When applying the effective Hamiltonian and the generalized factorization approximation to $B_c^- \to D_s^- D^0(\bar{D}^0)$ decay, one has to take the possible final state interactions into account. Since here both the final states i.e., $D_s^- D^0$ and $D_s^- \bar{D}^0$ are isospin 1/2 states, in general one can expect that both the amplitudes have

same strong FSI phases. Thus the amplitudes for these decay modes are given as

$$A(B_c^- \to D_s^- \bar{D}^0) = \frac{G_F}{\sqrt{2}}(V_{ub}V_{cs}^*)e^{i\delta}(a_1 X + a_2 Y)$$
$$A(B_c^- \to D_s^- D^0) = \frac{G_F}{\sqrt{2}}(V_{cb}V_{us}^*)e^{i\delta}(a_2 Y) \tag{3}$$

where δ is the strong FSI phase which is same for both the processes. X and Y are the factorized hadronic matrix elements

$$\begin{aligned} X &= \langle D_s^-|(\bar{s}c)|0\rangle\langle\bar{D}^0|(\bar{u}b)|B_c^-\rangle \\ Y &= \langle\bar{D}^0|(\bar{u}c)|0\rangle\langle D_s^-|(\bar{s}b)|B_c^-\rangle \end{aligned} \tag{4}$$

Since $V_{ub} = |V_{ub}|e^{-i\gamma}$ the weak phase difference between the amplitudes $A(B_c^- \to D_s^- \bar{D}^0)$ and $A(B_c^- \to D_s^- D^0)$ amounts to $-\gamma$ and there is no strong phase difference between them. As shown in Ref. [20] both these amplitudes are of the same order of magnitude:

$$r = \frac{|A(B_c^- \to D_s^- \bar{D}^0)|}{|A(B_c^- \to D_s^- D^0)|} = \mathcal{O}(1)\ . \tag{5}$$

Now let us write the amplitudes in a more generalized form i.e., in terms of their magnitudes, strong and weak phases

$$\begin{aligned} A(B_c^- \to D^0 D_s^-) &= A_1 e^{i\delta} \\ A(B_c^- \to \bar{D}^0 D_s^-) &= A_2 e^{-i\gamma} e^{i\delta} \end{aligned} \tag{6}$$

where A_1 and A_2 are the magnitudes of the corresponding amplitudes, δ is the strong phase and γ is the weak phase. These forms of the amplitudes give $r = A_2/A_1$. The amplitudes for the corresponding charge conjugate states are given as

$$\begin{aligned} A(B_c^+ \to \bar{D}^0 D_s^+) &= A_1 e^{i\delta} = A(B_c^- \to D^0 D_s^-) \\ A(B_c^+ \to D^0 D_s^+) &= A_2 e^{i\gamma} e^{i\delta} = e^{2i\gamma}\ A(B_c^- \to \bar{D}^0 D_s^-) \end{aligned} \tag{7}$$

The decay rates for the flavor specific states of D mesons are given as (disregarding the phase space factor)

$$\begin{aligned} \Gamma(B_c^- \to D^0 D_s^-) &= \Gamma(B_c^+ \to \bar{D}^0 D_s^+) = A_1^2 \\ \Gamma(B_c^- \to \bar{D}^0 D_s^-) &= \Gamma(B_c^+ \to D^0 D_s^+) = A_2^2 \end{aligned} \tag{8}$$

Now from Eqs. (6) and (7) we can write the amplitudes for the decay of $B_c^\pm$ into CP eigen state $D_+^0 (= (D^0 + \bar{D}^0)/\sqrt{2})$ and $D_s^\pm$ as

$$\sqrt{2}A(B_c^+ \to D_+^0 D_s^+) = A(B_c^+ \to D^0 D_s^+) + A(B_c^+ \to \bar{D}^0 D_s^+)$$
$$\sqrt{2}A(B_c^- \to D_+^0 D_s^-) = A(B_c^- \to D^0 D_s^-) + A(B_c^- \to \bar{D}^0 D_s^-) \qquad (9)$$

From the two expressions given in Eq. (9) one can construct two triangles with the common side $A(B_c^+ \to \bar{D}^0 D_s^+) = A(B_c^- \to D^0 D_s^-)$ and the angle (2γ) can be determined. This method is recently described by Fleischer and Wyler [20]. The advantage of this method is that here all sides of the two triangles are of comparable length giving rise to nonsquashed triangles.

However, here we proceed in a different manner analogous to Ref. [13]. We consider the decay of $B_c^\pm$ into both CP even $D_+^0((D^0+\bar{D}^0)/\sqrt{2})$ and CP odd $D_-^0((D^0-\bar{D}^0)/\sqrt{2})$ alongwith the accompanying $D_s^\pm$ meson i.e.$B_c^\pm \to D_{+,-}^0 D_s^\pm$, modes. The amplitudes for $B_c^\pm \to D_-^0 D_s^\pm$ can be written in the same form as Eq. (9) by changing the sign of second terms. Thus neglecting the small $D^0 - \bar{D}^0$ mixing we obtain the decay rates into CP eigen states of final D meson as (the common phase space factors are not taken into account)

$$\Gamma(B_c^\pm \to D_+^0 D_s^\pm) = \frac{1}{2}\left[A_1^2 + A_2^2 + 2A_1A_2\cos\gamma\right]$$
$$\Gamma(B_c^\pm \to D_-^0 D_s^\pm) = \frac{1}{2}\left[A_1^2 + A_2^2 - 2A_1A_2\cos\gamma\right] \qquad (10)$$

Now we define two charge-averaged ratios for the two CP eigenstates

$$\begin{aligned} R_i &= 2\,\frac{\Gamma(B_c^+ \to D_i^0 D_s^+) + \Gamma(B_c^- \to D_i^0 D_s^-)}{\Gamma(B_c^+ \to \bar{D}^0 D_s^+) + \Gamma(B_c^- \to D^0 D_s^-)} \\ &= 1 + r^2 \pm 2r\cos\gamma\ , \qquad \text{where} \qquad i = +,- \end{aligned} \qquad (11)$$

This equation can be written in a more generalized form as

$$R_{+,-} = \sin^2\gamma + (r \pm \cos\gamma)^2 \qquad (12)$$

from which one may get the constraint

$$\sin^2\gamma \leq R_{+,-}\ . \qquad (13)$$

The weak phase γ can be written in terms of R_+ and R_- from Eq. (11) as

$$\cos\gamma = \frac{1}{4}\frac{R_+ - R_-}{\sqrt{\frac{1}{2}(R_+ + R_-) - 1}}\ . \qquad (14)$$

Thus the unknown γ can be easily determined in terms of R_+ and R_-, of course with four fold quadrant ambiguities. The CP even (odd) state can be identified by its CP even (odd) decay products. For instance the states $K_s\pi^0$, $K_s\rho$, $K_s\omega$, $K_s\phi$ can be used to identify D_-^0 while $\pi^+\pi^-$, K^+K^- represent a D_+^0.

The advantage of these decay modes $B_c^\pm \to D_s^\pm\{D^0, \bar{D}^0, D_+^0, D_-^0\}$ is that there is no FSI strong phase difference in these decay modes since all these modes have only the isospin 1/2 final states. So these modes can in principle be considered as gold plated modes for the extraction of angle γ. As discussed by Fleischer and Wyler [20] one expects a huge number of B_c mesons, about 10^{10} untriggered B_c's per year of running and expects around 20 events per year at LHC for an overall efficiency of 10%. Hence it seems B_c system may well contribute to our understanding of CP violation.

3 γ from $B_c^\pm \to D_s^{*\pm} D^{*0}(\bar{D}^{*0})$

The angular distribution for $B_c^- \to D_s^{*-} D^{*0}(\bar{D}^{*0}) \to (D_s^- \gamma)(D^0(\bar{D}^0)\pi)$ which is same as $B_c^\pm \to V_1 V_2 \to (l^+ l^-)(P\pi)$ [22] is given in the linear polarization basis as [23]

$$\frac{1}{\Gamma}\frac{d^3\Gamma}{d\cos\theta \; d\cos\psi \; d\phi} = \frac{9}{32\pi}\Big[2|A_0|^2\cos^2\psi(1-\sin^2\theta\cos^2\phi)$$
$$+\sin^2\psi\left\{|A_\parallel|^2(1-\sin^2\theta\sin^2\phi)+|A_\perp|^2\sin^2\theta - \mathrm{Im}(A_\parallel^* A_\perp)\sin 2\theta\sin\phi\right\}$$
$$+\frac{1}{\sqrt{2}}\sin 2\psi\left\{\mathrm{Re}(A_0^* A_\parallel)\sin^2\theta \;\; \sin 2\phi + \mathrm{Im}(A_0^* A_\perp)\sin 2\theta\cos\phi\right\}\Big] \qquad (15)$$

where $A_\perp$ is the P wave decay amplitude, A_0 and $A_\parallel$ are two orthogonal combinations of the S and D wave amplitudes with the normalization $|A_\perp|^2 + |A_0|^2 + |A_\parallel|^2 = 1$. The angular distribution of $B_c^+ \to D_s^{*+} D^{*0}(\bar{D}^{*0})$ decay is similar to Eq. (15) and we shall use $\bar{A}$ to indicate the corresponding decay amplitudes. Each amplitude A_i has both CP conserving FSI phase δ_i and CP violating weak phase σ_i i.e. $A_i = |A_i| e^{i(\delta_i + \sigma_i)}$ while the corresponding amplitudes $\bar{A}_i$ are related to A_i as

$$\bar{A}_\perp = -|A_\perp| e^{i(\delta_\perp - \sigma_\perp)} , \quad \bar{A}_\parallel = |A_\parallel| e^{i(\delta_\parallel - \sigma_\parallel)} , \quad \text{and} \quad \bar{A}_0 = |A_0| e^{i(\delta_0 - \sigma_0)} . \qquad (16)$$

The rich kinematics of the vector-vector final state allows one to separate each of the six combinations of the linear polarization amplitudes of Eq. (15). The weight factors for the corresponding amplitudes can be determined as done in Ref. [17] for $B \to V_1 V_2 \to (P_1\pi)(P_2\pi)$ using the Fourier transform in ϕ and orthogonality of Legendre polynomial in $\cos\theta$ and $\cos\psi$. An observable can be determined from its weight factor W_i, given in table-1, using

$$O_i = \frac{32\pi}{9}\int d\cos\theta \; d\cos\psi \; d\phi \frac{W_i}{\Gamma}\frac{d^3\Gamma}{d\cos\theta \; d\cos\psi \; d\phi} . \qquad (17)$$

It should be noted that in this case the weight factors do not give identical results under the interchange of $\theta \leftrightarrow \psi$ as in the case of Ref. [17]. The

amplitudes for B_c^- decays for a given polarization state is given as

$$\begin{aligned} A^\lambda(B_c^- \to D_s^{*-} D^{*0}) &= |V_{cb}V_{us}^*|\ a_1^\lambda\ e^{\delta^\lambda} \\ A^\lambda(B_c^- \to D_s^{*-} \bar{D}^{*0}) &= |V_{ub}V_{cs}^*|\ a_2^\lambda\ e^{\delta^\lambda} e^{-i\gamma} \end{aligned} \tag{18}$$

Since both the final states $D_s^{*-}D^{*0}$ and $D_s^{*-}\bar{D}^{*0}$ have isospin 1/2, we have taken same strong FSI phase δ^λ for both the decay modes. The amplitudes for the corresponding charge conjugate modes can be obtained using Eq. (16), as

$$\begin{aligned} A^\lambda(B_c^+ \to D_s^{*+} \bar{D}^{*0}) &= x^\lambda |V_{cb}^*V_{us}|\ a_1^\lambda\ e^{i\delta^\lambda} \\ A^\lambda(B_c^+ \to D_s^{*+} D^{*0}) &= x^\lambda |V_{ub}^*V_{cs}|\ a_2^\lambda\ e^{i\delta^\lambda} e^{i\gamma} \end{aligned} \tag{19}$$

where $x^\lambda = -1$ for $\lambda =\perp$ and $+1$ for $\lambda =\|$ and 0. We now consider the decay of $D^{*0}/\bar{D}^{*0}$ into $D^0\pi^0/\bar{D}^0\pi^0$ with $D^0/\bar{D}^0$ meson further decaying to a common final state f. Generally for B decays f is chosen to be a Cabibbo allowed mode for $\bar{D}^0$ while it is doubly Cabibbo suppressed for D^0 because in that case the ratio of the amplitudes $|A(B^- \to \bar{D}^0K^-)|/|A(B^- \to D^0K^-)| \approx \mathcal{O}(0.1)$. However, in B_c decays the ratios of the amplitudes is $\mathcal{O}(1)$, so here we take two possible cases : one as done in B case i.e. $D^0 \to f$ is doubly Cabibbo suppressed while $\bar{D}^0 \to f$ is Cabibbo allowed and the other with f as a CP eigen state.

Let us first consider the first case where f can be taken as $K^+\pi^-$. Neglecting the negligible mixing effects in the $D^0 - \bar{D}^0$ system the amplitudes for the decays of $B_c^\pm$ to the final state f and its CP conjugate can be written as

$$\begin{aligned} A_f^\lambda &= A^\lambda(B_c^- \to [[f]_D\pi]_{D^*}D_s^{*-}) \\ &= \sqrt{B}e^{i\delta^\lambda}\left[|V_{ub}V_{cs}^*|\ a_2^\lambda\ e^{-i\gamma} + |V_{cb}V_{us}^*|\ R\ a_1^\lambda\ e^{i\Delta}\right] \\ \bar{A}_{\bar{f}}^\lambda &= A^\lambda(B_c^+ \to [[f]_D\pi]_{D^*}D_s^{*+}) \\ &= x^\lambda\sqrt{B}e^{i\delta^\lambda}\left[|V_{ub}^*V_{cs}|\ a_2^\lambda\ e^{i\gamma} + |V_{cb}^*V_{us}|\ R\ a_1^\lambda e^{i\Delta}\right] \\ A_{\bar{f}}^\lambda &= A^\lambda(B_c^- \to [[\bar{f}]_D\pi]_{D^*}D_s^{*-}) \\ &= \sqrt{B}e^{i\delta^\lambda}\left[|V_{ub}V_{cs}^*|\ R\ a_2^\lambda\ e^{-i\gamma}e^{i\Delta} + |V_{cb}V_{us}^*|\ a_1^\lambda\ \right] \\ \bar{A}_f^\lambda &= A^\lambda(B_c^+ \to [[f]_D\pi]_{D^*}D_s^{*+}) \\ &= x^\lambda\sqrt{B}e^{i\delta^\lambda}\left[|V_{ub}^*V_{cs}|\ R\ a_2^\lambda\ e^{i\gamma}e^{i\Delta} + |V_{cb}^*V_{us}|\ a_1^\lambda\ \right] \end{aligned} \tag{20}$$

where $[X]_M$ indicates that the state X is constructed to have invariant mass of M; $B = Br(\bar{D}^0 \to f)$, $R^2 = Br(\bar{D}^0 \to f)/Br(D^0 \to f)$ and Δ is the strong phase difference between $\bar{D}^0 \to f$ and $\bar{D}^0 \to \bar{f}$.

Thus the measurement of the angular distribution for each of the four modes provides us with a total of twentyfour observables, six for each mode. These observables can be extracted experimentally using the weight functions.

There are only thirteen unknowns in Eq.(20) : R, Δ, γ, $|V_{ub}|$ and three variables for each a_1^λ, a_2^λ and δ^λ. Thus γ in principle can be easily determined from these observables.

Now let us consider f to be a CP eigenstate i.e. $f = K^+K^-$ or $\pi^+\pi^-$ with CP eigenvalue $+1$. In this case the number of unknowns in Eq. (20) is further reduced because there is no relative strong phase difference between $\bar{D}^0 \to f$ and $\bar{D}^0 \to \bar{f}$ as f is a CP eigen state. So there will no longer be the strong phase difference factor ($e^{i\Delta}$) in the expressions for the amplitudes (20). Furthermore, since f is chosen to be a CP eigenstate R is also no longer an unknown. It can be related to $Br(\bar{D}^0 \to f)$ through the experimentally determined CP rate asymmetries $a_{CP}(f)$. Defining $a_{CP}(f)$ as

$$a_{CP}(f) = \frac{Br(\bar{D}^0 \to f) - Br(D^0 \to f)}{Br(\bar{D}^0 \to f) + Br(D^0 \to f)} \tag{21}$$

which gives

$$R = \frac{1 - a_{CP}(f)}{1 + a_{CP}(f)} \tag{22}$$

Since a_{CP} is very small i.e., $a_{CP}(K^+K^-) = 0.026 \pm 0.03$ and $a_{CP}(\pi^+\pi^-) = -0.05 \pm 0.08$ [24], R can be taken as approximately 1.

Thus we get rid of two more unknowns R and Δ if we consider the common final state f to be a CP eigenstate. In this case the total number of unknowns are eleven which would possibly be overdetermined with the twentyfour observables.

Now let us consider whether the decay modes $B_c^\pm \to D_s^{*\pm} D^{*0}(\bar{D}^{*0})$ can be used to observe CP violation or not. It is well known that to observe direct CP violation one would require two interferring amplitudes with different strong and weak phases. Thus the usual signature of CP violation i.e. the partial rate asymmetry

$$|A_f^\lambda|^2 - |\bar{A}_{\bar{f}}^\lambda|^2 = 4|V_{ub}^* V_{cs} V_{cb} V_{us}^*| \ R \ B \ a_1^\lambda \ a_2^\lambda \ \sin\Delta \sin\gamma \ , \tag{23}$$

can not be observed in case f is chosen to be a CP eigen state. It is therefore interesting to see if there are other observables in the angular distribution which can provide useful information about CP violation even if partial rate asymmetries are zero. It is clear that the coefficients $\hat{\alpha} = -\mathrm{Im}(A_\parallel^* A_\perp)$, $\hat{\gamma} = \mathrm{Im}(A_0^* A_\perp)$, the fourth and last terms in the expression for angular distribution (15), and similarly $\bar{\hat{\alpha}}$ and $\bar{\hat{\gamma}}$ for B_c^+ decay, contain information about CP violation, even for cases where the two amplitudes have no relative strong phase difference. From Eq. (20) we can obtain the following quantities which can measure CP violation as

$$\begin{aligned}&\mathrm{Im}\{(A^\perp A^{\rho *})_f + (\bar{A}^\perp \bar{A}^{\rho *})_{\bar{f}}\}\\ &= 2|V_{ub}^* V_{cs} V_{cb} V_{us}^*| RB \sin\gamma [a_1^\perp a_2^\rho \cos(\delta^\perp - \delta^\rho + \Delta) - a_2^\perp a_1^\rho \cos(\delta^\perp - \delta^\rho - \Delta)]\end{aligned}$$

$$\begin{aligned}
&\mathrm{Im}\{(A^{\perp}A^{\rho *})_{\bar f} + (\bar A^{\perp}\bar A^{\rho *})_{f}\} \\
&= 2|V_{ub}^{*}V_{cs}V_{cb}V_{us}^{*}|RB\sin\gamma[a_1^{\perp}a_2^{\rho}\cos(\delta^{\perp}-\delta^{\rho}-\Delta) - a_2^{\perp}a_1^{\rho}\cos(\delta^{\perp}-\delta^{\rho}+\Delta)] \\
&\mathrm{Im}\{(A^{\perp}A^{\rho *})_{f} + (\bar A^{\perp}\bar A^{\rho *})_{\bar f} + (A^{\perp}A^{\rho *})_{\bar f} + (\bar A^{\perp}\bar A^{\rho *})_{f}\} \\
&= 4|V_{ub}^{*}V_{cs}V_{cb}V_{us}^{*}|RB\sin\gamma\cos(\Delta)\cos(\delta^{\perp}-\delta^{\rho})[a_1^{\perp}a_2^{\rho} - a_2^{\perp}a_1^{\rho}]\ , \qquad (24)
\end{aligned}$$

where $\rho = \|$ or 0. When $\rho = \|$ we observe CP violation in $\hat\alpha$ parameter and $\rho = 0$ corresponds to $\hat\gamma$ asymmetries. These CP violating observables do not require FSI phase differences and are especially sensitive to CP violating weak phases. Thus unlike $B_c^{\pm} \to D_s^{\pm} D^0(\bar D^0)$ decays where direct CP violation can not be observed, whereas in this case one can observe the signatures of CP violation. However, the weak phase γ can be cleanly determined in both the cases.

Now let us make a crude estimate of the number of events required to measure γ by this method, assuming that about 10^{10} untriggered B_c's will be available at LHC per year of running. The observed number of events for each mode is [14]

$$N_{obs} = N_0 \times Br \times f \times \epsilon\ , \qquad (25)$$

where N_0, Br, f and ϵ are the total number of B_c events, branching ratio, observation fraction and detector efficiency, respectively. The particle $\bar D^0$ is seen in its flavor tagging $K^+\pi^-$ and $K^+\pi^-\pi^+\pi^-$ modes. Ref. [11] gives the list of visible fraction (f) and the detector efficiencies (ϵ) for the various final state particles. The decay width for the color suppressed mode $B_c^+ \to D_s^{*+}D^{*0}$ is estimated in Ref. [25] with value $\Gamma(B_c^+ \to D_s^{*+}D^{*0}) = a_2^2\ 0.564 \times 10^{-13}$ MeV. Using $a_2 = 0.23$, the branching ratio is found to be

$$Br(B_c^+ \to D_s^{*+}D^{*0}) = 2 \times 10^{-6}\ . \qquad (26)$$

For the decay mode $B_c^+ \to D_s^{*+}\bar D^{*0}$ we expect a branching ratio at 10^{-5} level.

If we assume the visible fraction (f) of the final D meson to be 11.5% and al overall efficiency of 10% [11] we get approximately 230 reconstructed events. This crude estimate indicates that the B_c system may well contribute to our understanding of CP violation.

4 Conclusion

In this paper we have discussed the determination of the angle γ from the pure tree nonleptonic B_c decay modes $B_c^{\pm} \to D_s^{\pm}D^0(\bar D^0)$ and $B_c^{\pm} \to D_s^{*\pm}D^{*0}(\bar D^{*0})$. For the former case we have followed the method of Gronau [13]. However, the decay modes $B_c^{\pm} \to D_s^{\pm}D^0(\bar D^0)$ are particularly interesting due to the following reasons. First they are free from the presence of the relative strong phases between $B_c^+ \to D_s^+ D^0$ and $B_c^+ \to D_s^+\bar D^0$ as they are both isospin 1/2 states. Secondly the two interferring amplitudes are of the same order of

magnitude i.e., the ratio of their amplitudes $r = |A(B_c^- \to D_s^- \bar{D}^0)|/|A(B_c^- \to D_s^- D^0)| \approx \mathcal{O}(1)$ whereas for the corresponding analog system $B^\pm \to DK^\pm$, r is $\mathcal{O}(0.1)$.

For the vector-vector final states we have followed the method of Ref [17]. Because of no relative strong phase between the two interferring amplitudes the number of unknowns are much less than the number of observables. Since the magnitudes of the amplitudes are of the same order here we have considered two possibilities for the common final state f $(D^{*0}/\bar{D}^{*0} \to D^0\pi^0/\bar{D}^0\pi^0 \to f\pi^0/f\pi^0)$. (i) f is a Cabibbo allowed mode for $\bar{D}^0$ and hence doubly Cabibbo suppressed for D^0. (ii) f is a CP eigen state. Due to the rich kinematics of vector-vector final states in this case one can observe alternative signature of CP violation even when there is zero partial rate asymmetry.

To summarize, it is possible to determine the weak phase γ from the measurement of the nonleptonic decay modes $B_c^\pm \to D_s^\pm D^0(\bar{D}^0)$ and the corresponding vector vector modes $B_c^\pm \to D_s^{*\pm} D^{*0}(\bar{D}^{*0})$ cleanly without any hadronic uncertainties (since they are free from the presence of relative strong phase difference and penguin pollutions). Further, in the vector-vector modes the determination is even cleaner as the number of observables is much more than the number of unknowns. Hence, these decay modes can in principle be considered as gold plated modes for the determination of the angle γ.

5 Acknowdgements

AKG would like to thank Council of Scientific and Industrial Research, Government of India, for financial support.

References

[1] *CP Violation* ed. C. Jarlskog (World Scientific, Singapore, 1989)

[2] *CP Violation* by I. I. Bigi and A. I. Sanda, Cambridge Monographs on Particle Physics, Nuclear Physics and Cosmology, 2000).

[3] N. Cabibbo, Phys. Rev. Lett. **10**, 531 (1963); M. Kobayashi and T. Maskawa, Prog. Theo. Phys. **49**, 652 (1973).

[4] A. J. Buras and R. Fleischer, in *Heavy Flavours II* eds. A. J. Buras and M. Lindner (World Scientific, Singapore, 1998) p. 65.

[5] *The BaBar Physics Book* eds. P. F. Harrison and H. R. Quinn (SLAC report 504, 1998).

[6] L. L. Chau and W. -Y. Keung, Phys. Rev. Lett. **53**, 1802 (1984); C. Jarlskog and R. Stora, Phys. Lett. **B 208**, 268 (1988).

[7] For a review see *B Decays* ed. S. Stone (World Scientific, Singapore 1994).

[8] I.I. Bigi and A. I. Sanda, Nucl. Phys. **B 193**, 85 (1991).

[9] M. Gronau and D. London, Phys. Rev. Lett **65**, 3381 (1990).

[10] M. Gronau and D. London, Phys. Lett. **B 253**, 483(1991): M. Gronau and D. Wyler, Phys. Lett. **B 265**, 172 (1991).

[11] I. Dunietz, Phys. Lett. B 270, 75 (1991).

[12] D. Atwood, I. Dunietz and A. Soni, Phys. Rev. Lett **78**, 3257 (1997); Phys. Rev. **D 63**, 036005 (2001).

[13] M. Gronau, Phys. Rev. **D 58**, 037301 (1998).

[14] J.-H. Jang and P. Ko, Phys. Rev **D 58**, 111302 (1998).

[15] A. F. Falk, A. A. Petrov, Phys. Rev. Lett. **85**, 252 (2000).

[16] M. Gronau, J. L. Rosner and D. London, Phys. Rev. Lett. **73**, 21 (1994); R. Fleischer, Phys. Lett. **B 365**, 399 (1996); M. Neubert and J. L. Rosner, Phys. Rev. Lett. **81**, 5076 (1998).

[17] N. Sinha and R. Sinha, Phys. Rev. Lett. **80**, 3706 (1998).

[18] N. G. Deshpande and C. O. Dib, Phys. Lett. **B 319**, 313 (1993).

[19] A. F. Falk, Y. Nir and A. A Petrov, Jour. of High Energy Phys. **9912**, 019 (1999) [hep-ph/9911369].

[20] R. Fleischer and D. Wyler, Phys. Rev. **D 62**, 057503 (2000).

[21] A. J. Buras and P. H. Weisz, Nucl. Phys. **B 333**, 66 (1990).

[22] M. Kramer., Phys. Rev. **D 45**, 193 (1992); Phys. Lett. **B 279**, 181 (1992).

[23] A. S. Dighe, I. Dunietz, H. J. Lipkin and J. L. Rosner, Phys. Lett. **B 369**, 144 (1996)

[24] D. E. Groom et al, Review of Particle Physics, The Euro. Phys. Jour. **C 15**, 1 (2000).

[25] C.-H. Chang and Y.-Q. Chen, Phys. Rev. **D 49**, 3399 (1994).

Table 1: The weight factors corresponding to the observables in the angular distribution for $B_c \to D_s^* D^{*0}$ decay modes.

Observable	Weight Factor
$\|A_0\|^2$	$\frac{9}{64\pi}\left(5\cos^2\psi - 1\right)$
$\|A_\parallel\|^2$	$\frac{9}{64\pi}\left(8\cos^2\phi - 5\sin^2\theta\right)$
$\|A_\perp\|^2$	$\frac{9}{32\pi}\left(2 - 5\cos^2\theta\right)$
$\mathrm{Im}(A_\parallel^* A_\perp)$	$\frac{15}{\pi^2}\cos^2\psi\cos\theta\sin\phi$
$\mathrm{Re}(A_0^* A_\parallel)$	$\frac{3}{\pi^2}\sqrt{2}\cos\psi\sin 2\phi$
$\mathrm{Re}(A_0^* A_\perp)$	$\frac{16}{\pi^3}\sqrt{2}\cos\psi\cos\theta\cos\phi$